J/ψ and Υ at high temperature

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Abstract

We use the screened Coulomb potential with r-dependent coupling constant and the non–perturbative quark–antiquark potential derived within the Field Correlator Method (FCM) to calculate J/ψ and Υ binding energies and melting temperatures in the deconfined phase of quark-gluon plasma.

keywords: Quark-gluon plasma, non–perturbative potential, heavy mesons
PACS: 12.38Lg, 14.20Lq, 25.75Mq

1 INTRODUCTION

Since 1986, the gold-plated signature of deconfinement was thought to be J/ψ suppression [1]. If Debye screening of the Coulomb potential above T_c is strong enough then J/ψ production in A+A collisions will be suppressed. Indeed, applying the Bargmann condition [2] for the screened Coulomb potential

$$V_C(r) = -\frac{4}{3}\cdot\frac{\alpha_s(r)}{r}\cdot e^{-m_d\, r}, \tag{1}$$

where m_d is the Debye mass, we obtain the simple estimate for the number of, say, $c\bar{c}$ S–wave bound states

$$n \leq \mu_c \int\limits_0^\infty |V_C(r)| r\, dr \;=\; \frac{4\alpha_s}{3}\cdot\frac{\mu_c}{m_d}\,, \tag{2}$$

where μ_c is the constituent mass of the c-quark and for the moment we neglect the r-dependence of α_s. Taking $\mu_c = 1.4$ GeV and $\alpha_s = 0.39$ we conclude that if $m_d \geq 0.7$ GeV, there is no J/ψ bound state. Parenthetically, we note that no light or strange mesons ($\mu \sim 300 - 500$ MeV) survive. But this is not the full story.

There is a significant change of views on physical properties and underlying dynamics of quark–gluon plasma (QGP), produced at RHIC, see *e.g.* [3] and references there in. Instead of behaving like a gas of free quasiparticles – quarks and gluons, the matter created in RHIC interacts much more strongly than originally expected. Also, the interaction deduced from lattice studies is strong enough to support $Q\overline{Q}$ bound states. It is more appropriate to describe the non-perturbative (NP) properties of the QCD phase close to T_c in terms remnants of the non–perturbative part of the QCD force rather than a strongly coupled Coulomb force.

In the QCD vacuum, the NP quark-antiquark potential is $V = \sigma r$. At $T \geq T_c$, $\sigma = 0$, but that does not mean that the NP potential disappears. In a recent paper [4] we calculated binding energies for the lowest $Q\overline{Q}$ and QQQ eigenstates ($Q = c, b$) above T_c using the NP $Q\overline{Q}$ potential derived in the Field Correlator Method (FCM) [5] and the screened Coulomb potential with the strong coupling constant $\alpha_s = 0.35$ in Eq. (1). In this talk we extend our analysis to the case of the running $\alpha_s(r)$.

2 The Field Correlator Method as applied to finite T

The NP $Q\overline{Q}$ potential can be studied through the modification of the correlator functions, which define the quadratic field correlators of the nonperturbative vaccuum fields

$$< tr\, F_{\mu\nu}(x)\Phi(x,0)F_{\lambda\sigma}(0) > = \mathcal{A}_{\mu\nu;\lambda\sigma}\, D(x) + \mathcal{B}_{\mu\nu;\lambda\sigma}\, D_1(x), \tag{3}$$

where $\mathcal{A}_{\mu\nu;\lambda\sigma}$ and $\mathcal{B}_{\mu\nu;\lambda\sigma}$ are the two covariant tensors constructed from $g_{\mu\nu}$ and $x_\mu x_\nu$, $\Phi(x,0)$ is the Schwinger parallel transporter, x Euclidian.

At $T \geq T_c$, one should distinguish the color electric correlators $D^E(x)$, $D_1^E(x)$ and color magnetic correlators $D^H(x)$, $D_1^H(x)$. Above T_c, the color electric correlator $D^E(x)$ that defines the string tension at $T = 0$ becomes zero [6] and, correspondingly, $\sigma^E = 0$. The color magnetic correlators $D^H(x)$ and $D_1^H(x)$ do not produce static quark–antiquark potentials, they only define the spatial string tension $\sigma_s = \sigma^H$ and the Debye mass $m_d \propto \sqrt{\sigma_s}$ that grows with T.

The main source of the NP static $Q\overline{Q}$ potential at $T \geq T_c$ originates from the color–electric correlator function $D_1^E(x)$

$$V_{np}(r,T) \;=\; \int\limits_0^{1/T} d\nu(1-\nu T) \int\limits_0^r \lambda d\lambda\, D_1^E(x). \tag{4}$$

In the confinement region the function $D_1^E(x)$ was calculated in [7]

$$D_1^E(x) \;=\; B\,\frac{\exp(-M_0\,x)}{x}, \tag{5}$$

where $B = 6\alpha_s^f \sigma_f M_0$, α_s^f being the freezing value of the strong coupling constant to be specified later, σ_f is the sting tension at $T = 0$, and the parameter M_0 has the meaning of the gluelump mass. In what follows we take $\sigma_f = 0.18$ GeV^2 and $M_0 = 1$ GeV. Above T_c the analytical form of D_1^E should stay unchanged at least up to $T \sim 2\,T_c$. The only change is

$B \to B(T) = \xi(T)B$, where the factor $\xi(T) = \left(1 - 0.36\frac{M_0}{B}\frac{T-T_c}{T_c}\right)$ is determined by lattice data [9]. Integrating over λ one obtains

$$V_{np}(r,T) = \frac{B(T)}{M_0}\int\limits_0^{1/T}(1-\nu T)\left(e^{-\nu M_0} - e^{-\sqrt{\nu^2+r^2}}\ M_0\right)d\nu \ = \ V(\infty,T) - V(r,T) \tag{6}$$

Note that $V(\infty, T_c) \approx 0.5$ GeV that agrees with lattice estimate for the free quark-antiquark energy.

In the framework of the FCM, the masses of heavy quarkonia are defined as

$$M_{Q\bar{Q}} \ = \ \frac{m_Q^2}{\mu_Q} \ + \ \mu_Q \ + \ E_0(m_Q, \mu_Q), \tag{7}$$

$E_0(m_Q, \mu_Q)$ is an eigenvalue of the Hamiltonian $H = H_0 + V_{np} + V_C$, m_Q are the bare quark masses, μ_Q are the auxiliary fields that are introduced to simplify the treatment of relativistic kinematics. The auxiliary fields are treated as c–number variational parameters to be found from the extremum condition imposed on $M_{Q\bar{Q}}$ in Eq. (7). Such an approach allows for a very transparent interpretation of axiliary fields as the constituent masses that appear due to the interaction. Once m_Q is fixed, the quarkonia spectrum is described. The dissociation points are defined as those temperature values for which the energy gap between $V(\infty, T)$ and E_0 disappears.

3 Coulomb potential

We use the perturbative screened Coulomb potential (1) with the r-dependent QCD coupling constant $\alpha_s(r,T)$. Note that in the entire regime of distances for which at $T = 0$ the heavy quark potential can be described well by QCD perturbation theory $\alpha_s(r,T)$ remains unaffected by temperature effects at least up to $T \leqslant 3\,T_c$ and agrees with the zero temperature running coupling $\alpha_s(r,0) = \alpha_s(r)$. For our purposes, we find it convenient to define the r–dependent coupling constant in terms of the $\mathbf{q}^2$ dependent constant $\alpha_B(\mathbf{q}^2)$ calculated in the background perturbation theory (BPTh) [8]

$$\alpha_s(r) \ = \ \frac{2}{\pi}\int\limits_0^{\infty} dq\ \frac{\sin\, qr}{q}\ \alpha_B(\mathbf{q}^2). \tag{8}$$

The formula for $\alpha_B(\mathbf{q}^2)$ is obtained by solving the two-loop renormalization group equation for the running coupling constant in QCD:

$$\alpha_B(\mathbf{q}^2) \ = \ \frac{4\pi}{\beta_0\, t}\left(1 - \frac{\beta_1}{{\beta_0}^2}\,\frac{\ln t}{t}\right), \quad t \ = \ \ln\,\frac{\mathbf{q}^2 + m_B^2}{\Lambda_V^2}, \tag{9}$$

where β_i are the coefficients of the QCD β-function.

In Eq. (9) the parameter $m_B \sim 1$ Gev has the meaning of the mass of the lowest hybrid excitation. The result can be viewed as arising from the interaction of a gluon with

Table 1: J/ψ above the deconfinement region. $V(\infty, T)$ is the continuum threshold (a constant shift in the potential). Units are GeV or GeV^{-1}.

T/T_c	$V(\infty)$	μ_b	$E_0 - V(\infty)$	r_0	$M_{J/\psi}$
1	0.445	1.443	-0.011	8.23	3.235
1.2	0.368	1.423	- 0.003	10.07	3.171

background vacuum fields. We employ the values $\Lambda_V = 0.36\,\mathrm{GeV}$, $m_B = 0.95\,\mathrm{GeV}$, which lie within the range determined in Ref. [10]. The result is consistent with the freezing of $\alpha_B(r)$ with a magnitude 0.563 (see Table 4 of Ref. [11]. The zero temperature potential with the above choice of the parameters gives a fairly good description of the quarkonium spectrum [10]. At finite temperature we utilize the information on m_d in Eq. (1) from Ref. [12]. For pure-gauge SU(3) theory ($T_c = 275$ MeV) m_d varies between 0.8 GeV and 1.4 GeV, when T varies between T_c and $2\,T_c$.

4 Results

The solutions for the binding energy for the $1S$ J/ψ and Υ states are shown in Tables 1, 2. In these Tables we present the constituent quark masses μ_Q for $c\bar{c}$ and $b\bar{b}$, the differences $\varepsilon_Q = E_0 - V_{Q\overline{Q}}(\infty)$, the mean squared radii $r_0 = \sqrt{\overline{r^2}}$, and the masses of the $Q\overline{Q}$ mesons. We employ $m_c = 1.4$ GeV and $m_b = 4.8$ GeV. Note that, as in the confinement region, the constituent masses μ_Q only slightly exceed bare quark masses m_Q that reflect smallness of the kinetic energies of heavy quarks.

At $T = T_c$ we obtain the weakly bound $c\bar{c}$ state that disappears at $T \sim 1.3\,T_c$. The charmonium masses lie in the interval 3.2 - 3.3 GeV, that agrees with the results of Ref. [9]. Note that immediately above T_c the mass of the $c\bar{c}$ state is about 0.2 GeV higher than that of J/ψ. As expected, the Υ state remains intact up to the larger temperatures, $T \sim 2.3\,T_c$, see Table 2. The masses of the L = 0 bottomonium lie in the interval 9.7–9.8 GeV, about 0.2–0.3 GeV higher than 9.460 GeV, the mass of $\Upsilon(1S)$ at $T = 0$. At $T = T_c$ the $b\bar{b}$ separation r_0 is 0.25 fm that compatible with $r_0 = 0.28$ fm at $T = 0$. The 1S bottomoniium undergo very little modification till $T \sim 2\,Tc$. The results agree with those found previously for a constant $\alpha_s = 0.35$ [4]. We also mention that the melting temperatures for Ω_c and Ω_b calculated in [4] practically coincide with those for J/ψ and Υ.

Our results for $1S(J/\psi)$ are qualitatively agree with those of Refs. [13], [14] based on phenomenological $Q\overline{Q}$ potentials identified with the free quark-antiquark energy measured on the lattice while our melting temperature for $1S(\Upsilon)$ is much smaller than $T \sim (4-6)\,T_c$ found in Ref. [14].

This work was supported in part by RFBR Grants # 08-02-00657, # 08-02-00677, # 09-02-00629 and by the grant for scientific schools # NSh.4961.2008.2.

Table 2: $1S$ $b\bar{b}$ state above the deconfinement region.

T/T_c	$V(\infty,T)$	μ_b	$E_0(T)-V(\infty,T)$	r_0	$M_{b\bar{b}}$
1	0.445	4.948	-0.255	1.39	9.796
1.3	0.332	4.922	-0.158	1.69	9.777
1.6	0.237	4.894	- 0.084	2.23	9.755
2.0	0.134	4.854	- 0.022	4.23	9.712
2.2	0.090	4.831	- 0.006	6.77	9.684
2.3	0.070	4.821	- 0.002	8.32	9.668

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